

Gravitational and mass distribution effects on stationary superwinds II. Extended dark matter haloes

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ABSTRACT

In this second part, we generalize the results of the previous paper. We present an analytic superwind solution considering extended gravitationally-interacting dark-matter and baryonic haloes. The incorporation of the latter is critical, since they can have a substantial effect on the hydrodynamics of superwinds generated by massive galaxies. Although the presence of extended and massive haloes does not change the limit for the closed-box enrichment of galaxies established in the first paper, they can trigger an earlier activation of the open-box enrichment scenario, since their gravitational potentials can contribute to the inhibition of the free superwind. Moreover, the incorporation of the extended haloes will also enhance the physical setting behind the superwind model, as we consider mass distributions with properties that emulate the results of recent simulations of Λ CDM haloes.

Key words: hydrodynamics, gravitation, galaxies: starburst, ISM: jets and outflows

1 INTRODUCTION

In the first part of this series of papers (Añorve-Zeferino & Corona-Galindo 2010; hereafter Paper I), we presented a simple, spherically-symmetric galactic superwind model considering non-uniform dynamical mass distributions with analogue energy and mass injection rates. Galaxies were modeled in terms of four parameters: a characteristic object radius, r_{sc} ; the effective energy deposition rate, \dot{E}_{eff} ; the effective mass deposition rate, \dot{M}_{eff} ; and a normalized spatial distribution, ρ_s . The latter defined the distributions of the dynamical mass and of the mass and energy injection rates within $r < r_{\text{sc}}$.

The spatial distribution was assumed to follow a truncated version of the Dehnen profile (1993), which allows to recover truncated versions of a plateau-like and the Hernquist (1990) and Jaffe (1983) profiles as particular cases. As an initial simplification, we considered only the dynamical mass contained within r_{sc} , which was assumed to account for most of the galaxy dark matter (DM) and baryonic mass (BM). This simplification allowed us to make a direct comparison between our analytic formulation and the numerical results of Silich et al. (2010), whom assumed a uniform distribution of the relevant galaxy parameters. Since in their model they also considered only the gravitational field of the central region, we were able to analytically reproduce their numerical results as a particular case. However, the previous assumption is a maximal extrapolation of the results of Persic, Salucci & Stel (1996) and Salucci & Persic (1997), whom found that in *some* cases the presence of dark matter begins to be important well within the galaxy

optical radius, R_{opt} . They estimated that from the galaxy centres up to R_{opt} , the fraction of DM goes from 0% up to 30%-70% (see also Salucci et al. 2007).

Thus, the initial supposition, as well as the assumption of spherical symmetry, can only be adequately interpreted in the context of a zeroth-order approximation for evaluating the effect of the gravitational field of ellipsoidal galaxies on the inner superwind solution ($r \leq r_{\text{sc}}$). Similarly, the zeroth-order distortion of the external superwind hydrodynamical profiles ($r > r_{\text{sc}}$) is only adequately predicted for galaxies with either low masses or very diffuse DM and BM haloes.

For disc galaxies with large masses, the aforementioned simplification certainly does not hold, as the observed flat rotation curves of the extended discs *require* a significant amount of dark and baryonic matter outside of the bulge ($r_{\text{sc}} \sim r_{\text{bulge}}$). Needless to say, the presence of discs will also produce collimated flows. Furthermore, extended and fairly massive haloes can also have important repercussions on the superwind solution for the case of ellipsoidal galaxies ($r_{\text{sc}} \sim r_{\text{nucleus}}$), according to the properties of DM haloes derived by Persic et al. (1996) and Salucci & Persic (1997).

Nevertheless, the zeroth-order superwind solution that we presented in Paper I has the advantage of being analytic. Thus, it can be used to construct the solution for the case in which departures from spherical symmetry are important. In order to do this adequately for a wide range of galaxy masses, we need first to incorporate the effect of the external DM and BM halo under the assumption of spherical symmetry.

In this work, we obtain such a solution considering gravitationally-interacting external haloes (Section 2). We consider

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halo distributions that emulate the results of recent cosmological simulations (Section 3). Later, in Section 4, we obtain new limits for the open-box enrichment scenario and the existence of accelerating superwind solutions (see Paper I). Finally, in Section 5, we evaluate the impact of the halo gravitational potential on the superwind hydrodynamical profiles. The conclusions are presented in Section 6.

2 AN ANALYTIC SOLUTION FOR THE FREE SUPERWIND INCLUDING EXTENDED HALOES

When the haloes are included, the equation of conservation of energy outside of the galaxy characteristic radius, equation (9) in Paper I, transforms into

$$\frac{1}{r^2} \frac{d}{dr} \left[\rho u r^2 \left(\frac{1}{2} u^2 + (\eta + 1) \frac{P}{\rho} \right) \right] = -\rho u (\nabla \phi + \nabla \phi_h). \quad (1)$$

Above, the hydrodynamical variables are represented by their usual symbols, η is the polytropic index and $-\nabla \phi = -GM_{\text{DM}}/r^2$, where M_{DM} is the total dynamical mass within r_{sc} . Similarly, $-\nabla \phi_h = -GM_h(r)/r^2$, where $M_h(r)$ is the cumulative dynamical mass (i.e. DM+BM) of the external halo, which has a total mass M_H . We will allow the profile of the external halo to be defined either as a continuation of the internal profile or as a centrally truncated profile with different characteristics.

The integration of equation (1) yields a Bernoulli-like equation

$$\frac{1}{2} u^2 + (\eta + 1) \frac{P}{\rho} = \frac{1}{2} V_g^2 - \tau \phi - \phi_h, \quad (2)$$

where $\phi = -GM_{\text{DM}}/r$, $\phi_h(r)$ is the gravitational potential at $r > r_{\text{sc}}$ associated to the non-truncated version of the external halo, $\tau = 1 - M_H/M_{\text{DM}}$ accounts for truncation effects, and V_g is the asymptotic terminal speed, which is given by

$$\frac{1}{2} V_g^2 = \frac{1}{2} V_\infty^2 - \frac{\tau G M_{\text{DM}}}{r_{\text{sc}}} + \phi_h(r_{\text{sc}}) - \frac{(1 + \frac{1}{A})}{(5 - 2\alpha)} \frac{G M_{\text{DM}}}{r_{\text{sc}}}. \quad (3)$$

In the last equation, V_∞ is the effective terminal speed due to the thermalization of SNe ejecta and individual stellar winds inside of the galaxy, α determines the steepness of the central dynamical mass distribution ($r < r_{\text{sc}}$) and A its concentration [see equation (11) in Paper I]. Note that $\tau = 0$ implies an uninterrupted, *continuous* gravitational potential. Similarly, $\tau < 0$ implies a centrally truncated external halo with a mass larger than M_{DM} , and $0 < \tau < 1$ implies the opposite. When $\tau = 1$ there is no external halo, and thus ϕ_h is identically zero.

As in Paper I, we will work in terms of dimensionless variables. For the present case they are:

$$R = \frac{r}{r_{\text{sc}}}, \quad U = \frac{u^2}{V_g^2}, \quad \text{and} \quad \Phi = -\tau \frac{V_{\text{eg}}}{R} + \Phi_h(R); \quad (4)$$

where V_{eg} is given by

$$V_{\text{eg}} = \frac{v_e^2}{V_g^2} = \frac{2G M_{\text{DM}}}{r_{\text{sc}} V_g^2}, \quad (5)$$

and $\Phi_h(R)$ is $\phi_h(r)$ written in terms of R and normalized to $V_g^2/2$. The conservation laws can now be reduced to the *same* governing differential equation than in Paper I, see its equation (43).

Thus, within the theoretical framework developed in Paper I, it is very easy to prove that the transonic free superwind solution is given by

$$R = D U^{-1/4} \left[1 - U + \tau \frac{V_{\text{eg}}}{R} - \Phi_h(R) \right]^{-\eta/2}, \quad (6)$$

with

$$D = \left(\frac{1}{2\eta + 1} \right)^{1/4} \left\{ \frac{2\eta [1 + \tau V_{\text{eg}} - \Phi_h(1)]}{2\eta + 1} \right\}^{\eta/2}. \quad (7)$$

Again, as in our previous work, we will give preference to the parametric version of the solution:

$$x[x - \tau V_{\text{eg}} + R \Phi_h(R)]^{\frac{2\eta-3}{4}} = D_0 (y+1)^{-\frac{1}{2\eta+1}} (2\eta - y)^{-\frac{2\eta}{2\eta+1}}, \quad (8)$$

where y is a parameter that varies between 0 and 2η and

$$x = R + \tau V_{\text{eg}} - R \Phi_h(R), \quad (9)$$

$$y = \left[\frac{(2\eta + 1)U}{1 + \tau \frac{V_{\text{eg}}}{R} - \Phi_h(R)} - 1 \right], \quad (10)$$

and

$$D_0 = [1 + \tau V_{\text{eg}} - \Phi_h(1)] (2\eta)^{\frac{2\eta}{2\eta+1}}. \quad (11)$$

To obtain the hydrodynamical profiles, one just needs to follow the algorithm presented at the end of section 3.3 in Paper I. An advantage of the parametric solution is that it allows to work with just functions of R in the first two critical steps, related to equations (8) and (9). On the other hand, equation (6) involves both R and U . For $\eta = 3/2$ (equivalent to the case $\gamma = 5/3$, where γ is the adiabatic index) there is no need for a numerical root finder in the first step of our algorithm. In the second step however, its use will be most likely unavoidable, as the particular form of the assumed gravitational potential (i.e. of the external halo profile) is involved.

In Section 4, we will give the limit above which the stationary solution is disrupted in the external zone ($r > r_{\text{sc}}$) and the necessary condition for an accelerating stationary superwind solution. In order to do this, we will specify first the normalized potential Φ_h in the next section.

3 THE EXTENDED HALO PROFILES

How are the DM and BM distributed¹ outside of the galaxy characteristic radius? Since we have permitted centrally truncated profiles for the external halo, theoretically, we can choose practically any of the usually assumed distributions; e.g. a NFW profile, Navarro, Frenk & White 1997; a generalized NFW profile, Moore et al. 1999; a self-similar profile, Yoshikawa & Suto 1999; an isothermal profile, and so on. Given that the most commonly used profiles depend on at least two parameters, and given also the additional freedom introduced by our truncated halo scheme; there is a vast number of profiles and parameters that can give reasonable agreement with observational studies and with the predictions of cosmological simulations.

We will try to rely on physical insight for selecting the external halo profile that we will use in our model. Recent cosmological simulations carried out by Abadi et al. (2010) predict that dark matter haloes always contract as a result of galaxy formation. They also found that the contraction effect is substantially less pronounced than predicted by the adiabatic contraction model (Blumenthal et al. 1986). On similar grounds, according to the high-resolution N-body cosmological simulations of Λ CDM haloes carried out by

¹ As in Paper I, we will assume that their distributions are analogue.

Navarro et al. (2010), the departures from similarity in the velocity dispersion and density profiles correlate in such a way, that a power law for the spherically averaged pseudo-phase-space density is preserved, $\rho/\sigma^3 \propto r^{-1.875}$. They remarked that the index of the previous power law is identical to that of a Bertschinger's similarity solution for self-similar infall onto a point mass (in an Einstein-de Sitter Universe). They conclude that Λ CDM haloes are not strictly universal, but that the departure from similarity previously mentioned may be a fundamental structural property.

Bearing in mind the results described above, we deduce that the cases $\tau < 0$ and $0 < \tau < 1$ correspond to artificial mathematically induced constraints that make continuous the potential at $r = r_{sc}$ for arbitrarily-chosen external-halo profiles [see equation (2)]. The case $\tau = 1$ is physical, but corresponds to the zeroth-order approximation regarded as inadequate for some galaxies in Section 1. Thus, the case $\tau = 0$ is of special interest, as it implies an unforced continuity of the gravitational potential. It turns out that adequately chosen truncated Dehnen profiles satisfy naturally the latter condition.

For $r < r_{sc}$, the cumulative dynamical mass corresponding to a truncated Dehnen profile² is given by equation (13) in Paper I:

$$M(r) = M_{DM}(1+A)^{3-\alpha} \left(\frac{R}{R+A} \right)^{3-\alpha}. \quad (12)$$

We will also assume a truncated Dehnen profile for the external halo, but we will demand a cumulative mass of the form:

$$M_h(r) = M_{DM}(1+A_1)^{3-\alpha_1} \left(\frac{R}{R+A_1} \right)^{3-\alpha_1}. \quad (13)$$

At $R = 1$ we have that $M(1) = M_h(1)$. Note that for this, we do not require $A_1 = A$ nor $\alpha_1 = \alpha$. The last property can be interpreted in terms of a contraction of an initial spatial configuration of DM and BM with concentration A_1 and steepness α_1 which produced a new configuration with concentration A and steepness α for $r < r_{sc}$, or well, vice-versa, if other processes were involved (v.gr. angular momentum). On the other hand, a trivial but important relationship can be obtained from the condition $M(1) = M_h(1)$ by separating the baryonic and dark matter components:

$$r_{sc}[M_{bar} + M_{dark}] = r_{sc}[M_{bar} + M_{dark}]_1. \quad (14)$$

This could be interpreted as an integral equivalent of the equation for adiabatic collapse derived by Blumenthal et al. (1986). Additionally, given that the radial velocity dispersion associated to the Dehnen profile goes as $\sigma \sim r^{\alpha/2}$ when $r \rightarrow 0$, we are able to recover the index of the Bertschinger's power law near the centre of the galaxy when $\alpha = 3/4$. However, Navarro et al. (2010) obtained the index from radial averaging, which implies that α can adopt values within a wider range.

Note that in turn, the previous configurations could be interpreted as the result of the contraction of an unperturbed configuration away from the galaxy. This is equivalent to saying that a galaxy formed from the perturbation of an initial state (A_0, α_0) , and that after certain time, the perturbation bifurcated and produced two inner contracted states characterized by (A, α) and (A_1, α_1) . The first state characterizes the inner regions of the galaxy, $r < r_{sc}$. Then, the characteristic radius r_{sc} can be taken either as the radius of a galaxy nucleus or of a bulge. The second state characterizes the outer portions of the galaxy (e.g. a disc + DM). This

is in agreement with the aforementioned cosmological simulations, and it implies that galaxies carved out gravitational potential holes when they formed, and that they correspond to local depressions of an otherwise smoother gravitational potential.

Here, we are just interested in the superwind solution, so, in order to keep things simple, we will just consider the states (A, α) and (A_1, α_1) , i.e. we will ignore the depression of the reference gravitational potential (A_0, α_0) . The price that we will pay for this, as well as for the implicit analogue distribution of the baryonic and DM components assumed in our scheme, is that instead of (almost) 'perfectly' flat rotation curves up to 15 times the optical radius (Persic et al. 1996, Salucci & Persic 1997), the rotation curves will show some downwards skewness at large radii. They are however well above the curves corresponding to keplerian rotation of the baryonic mass. Evenmore, the behaviour of the associated rotation curves away from r_{sc} is consistent with that of the universal rotation curves derived by Salucci et al. (2007) for spiral galaxies. Anyway, for our purposes, the behaviour at large radii is not that important, as the thermalization driven superwind solution is valid only close to the galaxy³ (see e.g. Strickland & Heckman 2009). So, we will proceed to give the expression corresponding to the external gravitational potential.

By taking the limit $R \rightarrow \infty$ in equation (13), one finds that the total dynamical mass is given by $M_t = M_{DM}(1+A_1)^{3-\alpha_1}$. The expression of the associated gravitational potential for $0 \leq \alpha \leq 1$ is then similar to that given by equation (2) in Dehnen (1993):

$$\Phi_h(R) = -\frac{V_{eg}(1+A_1)^{3-\alpha_1}}{(2-\alpha_1)A_1} \left[1 - \left(\frac{R}{R+A_1} \right)^{2-\alpha_1} \right]. \quad (15)$$

With this, we can establish new approximated thresholds for the open-box enrichment scenario and for accelerating superwind solutions.

4 THRESHOLDS FOR OPEN-BOX ENRICHMENT AND ACCELERATING SUPERWIND SOLUTIONS

When the effect of the external halo is considered, the asymptotic terminal speed is given by

$$V_g = \left[1 - \left(\frac{1+\frac{1}{A}}{5-2\alpha} \right) V_e + \frac{V_g^2}{V_\infty^2} \Phi_h(1) \right]^{1/2} V_\infty. \quad (16)$$

The flow enters into non-stationary regimes (inpouring or outpouring, see Paper I) when

$$\left(\frac{1+\frac{1}{A}}{5-2\alpha} \right) V_e - \frac{V_g^2}{V_\infty^2} \Phi_h(1) \geq 1. \quad (17)$$

When the above inequality holds, the galaxy can eventually enter into an open-box enrichment scenario. Otherwise, we will have fully stationary solutions, unless radiative cooling or self-gravitation inhibit the stationary solution.

Fully stationary superwinds have accelerating velocity profiles when

$$-(2\eta+1) \frac{V_g^2}{V_\infty^2} \Phi_h(1) + 2\eta \left(\frac{1+\frac{1}{A}}{5-2\alpha} \right) V_e \leq 2\eta, \quad (18)$$

² See also equation (3) in Dehnen (1993).

³ This implies that the effect of the 'real' Φ_h can be emulated there by giving adequate values to A_1 and α_1 .

otherwise, they have decelerating velocity profiles. When the equality holds in the above relation, we have an almost constant external velocity profile with characteristic velocity

$$V_g = (2\eta)^{-1/2} [-2\phi_h(r_{sc})]^{1/2} = \eta^{-1/2} v_{rot} \left\{ \frac{(1+A_1)^{3-\alpha_1}}{(2-\alpha_1)A_1} \left[1 - \left(\frac{1}{1+A_1} \right)^{2-\alpha_1} \right] \right\}^{1/2}, \quad (19)$$

where v_{rot} is the rotation speed at r_{sc} .

5 EFFECT ON THE HYDRODYNAMICS

In our model, the dynamical mass (M_{DM}) contained within a bulge or galaxy nucleus experiencing a starburst episode is related to the concentration parameter and steepness of the external halo, and to the total dark matter and baryonic mass:

$$M_{DM} = \frac{M_t}{(1+A_1)^{3-\alpha_1}}. \quad (20)$$

The dynamical mass contained in the external halo ($r > r_{sc}$) is

$$M_H = M_t \left[1 - \frac{1}{(1+A_1)^{3-\alpha_1}} \right]. \quad (21)$$

Similarly, the dynamical mass contained up to an external characteristic radius (normalized to r_{sc}), R_D , is

$$M_D = M_t \left(\frac{R_D}{R_D + A_1} \right)^{3-\alpha_1}. \quad (22)$$

The radius R_D can be associated to the 'disc' radius of the BM or well to the BM+DM virial radius. So, all the relevant galaxy parameters are correlated, in a similar fashion as in the work of Salucci et al. (2007). Nevertheless, we emphasize that the relationship between the parameters is alike but of course not the same, since here we constructed our theoretical model only following the results of the simulations of Abadi et al. (2010) and Navarro et al. (2010).

We will proceed to discuss the effect of the extended haloes on the hydrodynamics. In order to do this, we consider the hydrodynamical models presented in Table 1 and Table 2. The first table gives the inner parameters for three galaxies with different characteristics. The second table gives the properties of their external haloes. The groundwork for the discussion will be the premise that the spherical symmetric superwind solution is a zeroth-order approximation to the aspherical case. We will consider again a reference effective terminal speed of 2500 km s^{-1} for the case of null mass-loading, fully efficient thermalization, and total participation within the starburst volume, i.e. for $\epsilon = \beta = \zeta = 1$. For models 2 and 3, the SFRs were obtained from formula (1) in Rupke, Veilleux & Sanders (2005a) and formula (28) in Paper I, i.e. we considered SFRs that are consistent with the typical observed luminosities for the object types, and that in parameter space, place the objects below the threshold for catastrophic cooling. We find that the predicted temperature profile is barely modified by the presence of the extended haloes. However, drastic changes are produced in the velocity profile.

The first model is an extended version of model 5 in Paper I, and corresponds to a synthetic isolated dwarf elliptical galaxy that tries to emulate the characteristics of the most massive outlier of the mass-metallicity relationship detected by Peeples, Pogge & Stanek (2008, see also Paper I). We assumed that the galaxy formed

by a contraction of $\sim 40\%$ of an initially unperturbed subhalo of DM and BM which had $\sim 70\%$ of its total mass located within $r \sim 3r_{sc}$, so we used $A_1 = 0.5$, $\alpha_1 = 3/4$ and $R_D = 1$. The latter is equivalent to saying that in this case there is no disc, i.e. we only have a galaxy nucleus. The internal dynamical mass distribution follows a plateau-like profile, which implies that some mechanism—perhaps the action of early powerful superwinds associated to a more extended and powerful starburst episode (see Governato et al. 2010) or internal dynamical processes—has also transformed the initial mass configuration. In this model, starburst activity still persists near the galaxy centre, but with a high concentration. We assumed a low thermalization efficiency, which implies a small number of massive stars and SNe within the characteristic concentration radius, $A = 0.1$. The justification for this is that the SFR is low, and that although small, the concentration radius is much larger than the typical radius of a massive star, i.e. the filling factor is low. Similarly, because of the small number of massive stars, just a small incorporation of mass is necessary to produce a heavily mass-loaded superwind. In this model, the presence of the extended halo suppresses the free superwind solution and the galaxy experiences an open-box enrichment [see equation (17)] by keeping the metals processed by the few massive stars still present near the galaxy centre. This will require however an already gas-poor galaxy at the moment at which the pollution occurred (Peeples, Pogge & Stanek, 2008). As suggested above, the required low mass fraction could have been produced by the action of early superwinds associated to previous and more powerful starburst activity. This is consistent with the views of Peeples et al. (2008), which regarded their sample of outliers as transitional galaxies in their way to becoming typically isolated dE and dSph galaxies, but with a high metallicity. The suppression of the free superwind solution is practically insensitive to the value of $0 \leq \alpha_1 \leq 1$, which indicates that the enrichment is produced by the physical conditions within A and the initial concentration of the unperturbed subhalo from which the galaxy formed.

The second model considers the synthetic and very massive blue compact dwarf galaxy modeled in Paper I. Here, we add an extended disc to the model in order to 'transform' the galaxy into a luminous infrared one⁴ (LIRG, $L_{IR} \sim 10^{11} L_\odot$). LIRGs and ULIRGs may be the end result of the merging of two moderate-size spiral galaxies and display traces of convergence to an elliptical morphology (Sanders & Mirabel 1996; Rupke et al. 2005a). We will model a LIRG that still exhibit evidence of heavily warped and thick discs, displaying a morphology perhaps similar to that of the central component of Arp 299 (Sargent & Scoville 1991; Heckman et al. 1999; Hibbard & Yun, 1999; Hu et al. 2004), but with just one nucleus. We assume that such features extend to up to 5 times the radius of the merger nucleus; thus, $R_D = 5$. The assumed mass and extension are consistent with CO emission observations of (U)LIRGs (Lonsdale, Farrah & Smith 2006 and references therein). We further assume that the merging process has similarly transformed the steepnesses of the internal and external mass profiles of the interacting galaxies unperturbed haloes, such that $\alpha = \alpha_1 = 3/4$. We adopt the value $A_1 = 1$ since it produces interesting proportions. In such a case, $\sim 66\%$ of the BM

⁴ N.B. As LIRGs and ULIRGs, BCDs may be the result of mergers, although generally they have lower masses, given that they mostly form from the merging of dwarf galaxies. Nevertheless, on a higher end, luminous blue compact galaxies can have dynamical masses of up to $\sim 10^{12} M_\odot$ (Garland et al. 2004, Pisano et al. 2010).

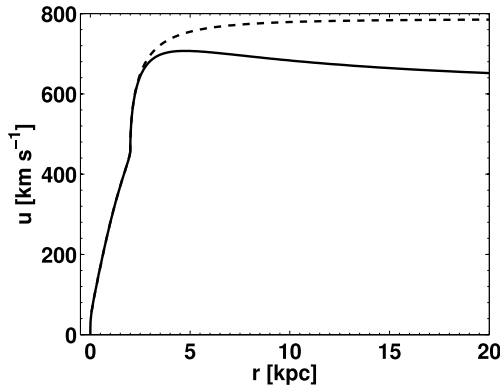


Figure 1. Superwind velocity profile for model 2 (solid line). The dashed-line represents the profile that would result if the external halo were neglected.

and DM of both galaxies is contained within the warped discs characteristic radius and about $\sim 30\%$ of this fraction resides within the merger nucleus (that is $\sim 20\%$ of the total mass). As in the original model, starburst activity is present in the nucleus with a somewhat high concentration ($A = 0.4$), the thermalization efficiency is 0.5 and mass loading is important, $\beta = 3$ (see Heckman et al. 1999). In this model, the gravitational field of the external halo transforms the accelerating superwind solution associated to the original model into a bounded decelerating one (Fig. 1). This effect occurs because now we have a more massive galaxy. The produced deceleration will enhance the observable properties of the superwind because of a proportional density increment. However, an even larger total mass could inhibit the superwind solution. As mentioned in Paper I, this is consistent with the superwind scaling properties found by Rupke, Veilleux & Sanders (2005b), whom reported an initial increment of the superwind observable properties with galaxy mass and a posterior flattening with the same.

Rupke et al. (2005b) also reported a flattening of the superwind observable properties at high SFR. In principle, the normalized free superwind solution (Section 2) is insensitive to the SFR (provided that it could be considered constant during a relatively large time interval), as it just depends on the effective and asymptotic terminal speeds. However, high SFRs will intensify the effect of radiative cooling, as more mass will be injected per unit time and volume, and thus, the stationary solution could also be radiatively inhibited. We have properly addressed this issue in section 2.1.2 of Paper I.

As an extreme example of the effect of the nominal value of the galaxy mass, we model a massive and ‘rare’ radio galaxy with a very extended halo (see e.g. Genzel et al. 2003). We consider a galaxy with a dynamical mass of $1 \times 10^{11} M_{\odot}$ within its nucleus of $r_{sc} \sim 2$ kpc. A mildly concentrated starburst ($A = 0.5$) is present in the nucleus, which has a cuspy dynamical mass distribution ($\alpha = 1$). We consider that the steepnesses of the inner region and the halo are the same and that the total mass of the galaxy is $M_t = 4 \times 10^{11} M_{\odot}$. This requires that $A_1 = 1$. This implies that the half-mass radius is $r \sim 2.5r_{sc}$ and that $\sim 80\%$ of the total mass is contained within $r \sim 8r_{sc}$. In this model a high deceleration of the superwind is produced, and the flow is unstable to small variations of the effective terminal speed (thermalization efficiency), as shown in Fig 2. As a consequence, the flow could eventually enter into the outpouring or even the inpouring regime. On the other hand, if instead of a continuous steepness, we consider that the typical cuspy halo profile (with slope $\alpha = 1$) resulted from the con-

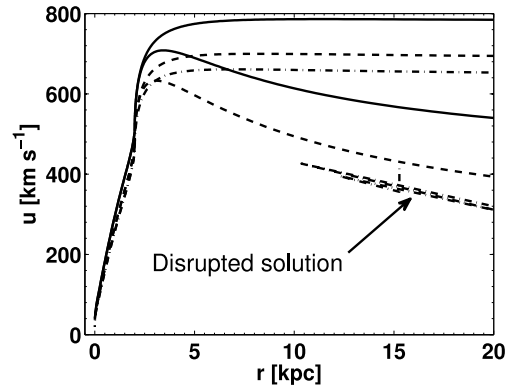


Figure 2. Velocity profiles for model 3. The solid lines corresponds to the parameters showed in Table 1. For this parameters, $V_{\infty} \approx 1208 \text{ km s}^{-1}$. The lower (upper) solid line (does not) consider(s) the presence of the external halo. Similarly, the dashed and dash-dotted lines correspond to $V_{\infty} = 1150 \text{ km s}^{-1}$ and $V_{\infty} = 1125 \text{ km s}^{-1}$, respectively. For the latter case, the stationary free superwind solution does not exist.

traction of a smoother one, say with $\alpha_1 = 3/4$ and $A_1 = 1$, the free superwind solution would be inhibited and the galaxy could enrich itself with is produced metals in an open-box scenario. This would occur because in the second case, the total mass is slightly larger, $M_t = 4.75 \times 10^{11} M_{\odot}$. The cumulative dynamical masses of the two assumed external profiles are very similar, their ratio varies from a value of 1 at r_{sc} (they are identical as they must), up to a value ~ 0.86 at $r = 10r_{sc}$; nevertheless, such a small variation is enough to suppress the stationary superwind solution. This reflects the fact that at the limit of large galaxy masses, galaxies will retain most of their metals, as expected.

6 SUMMARY AND CONCLUSIONS

Here we have presented an analytic free superwind model that incorporates the effect of extended DM and BM haloes. We find that the gravitational field of the extended haloes associated to massive galaxies can drastically alter the free superwind velocity profile and enhance its observable properties. We also find that massive haloes can also contribute to the inhibition of the superwind solution.

In our model, the galaxy total mass (BM+DM), the mass contained within a bulge or galaxy nucleus (defined by the characteristic radius r_{sc}), the mass up to the disc characteristic radius, and the steepness and concentration of the external halo, are all correlated. Since the correlations are nonlinear, deviations from galaxy to galaxy are permitted, see Tables 1 and 2. Oppositely, there is no correlation between the above parameters and the concentration and steepness of the mass distribution for $r < r_{sc}$. This is consistent with the results of the cosmological simulations carried out by Abadi et al. (2010) Navarro et al (2010), in the sense that haloes are not strictly universal. This should be expected, as we based our model in the ‘structural contraction’ property derived from their simulations. On the other hand, in their extensive work, Salucci et al. (2007) found that the previous parameters were correlated for spiral galaxies, and proposed universal rotation curves assuming a Burkert (1995) profile for the DM distribution. Our theoretical work diverges from theirs in that we considered additionally the mentioned inner concentration and steepness, which traces starburst episodes. Such a consideration discards the possibility of universal halo profiles and rotation curves, since in general they will

Table 1. Reference hydrodynamical models. Galaxy parameters for $r < r_{\text{sc}}$.

Model	Type	α	A	r_{sc} (kpc)	M_{DM} ($\times 10^8 M_{\odot}$)	SFR $M_{\odot} \text{ yr}^{-1}$	β	ϵ	ζ	V_{∞} km s^{-1}	V_e	Regime (No halo)
		(a)	(b)	(c)	(d)	(e)	(f)	(g)	(h)	(i)	(j)	(k)
1	dE	0	0.1	1	100	0.1	4	0.2	1	560	0.2740	borderline
2	(L)BCD/LIRG	0.75	0.4	2	500	~ 40	3	0.5	1	1021	0.2060	accelerating
3	Radio	1	0.5	2	1000	~ 200	3	0.7	1	1208	0.2946	accelerating

Superwind hydrodynamical models. Table headers: (a) steepness parameter, (b) concentration parameter, (c) radius, (d) dynamical mass, (e) star formation rate, (f) mass loading factor, (g) thermalization efficiency, (h) participation factor (i) effective terminal speed, (j) squared ratio of the escape velocity to the effective terminal speed, and (k) flow regime when the external halo is neglected.

Table 2. Reference hydrodynamical models. External halo parameters.

Model	Type	α_1	A_1	r_{D} (kpc)	M_{D} ($\times 10^8 M_{\odot}$)	M_{t} ($\times 10^8 M_{\odot}$)	Regime
		(a)	(b)	(c)	(d)	(e)	(f)
1	dE	0.75	0.5	1	M_{DM}	$\sim 2.5 M_{\text{DM}}$	open-box enrichment
2	LIRG	0.75	1	5	$\sim 3.16 M_{\text{DM}}$	$\sim 4.76 M_{\text{DM}}$	decelerating
3	Radio	1 (3/4)	1	$R_{\text{hm}} = 2.5$	$2 M_{\text{DM}}$	$4 M_{\text{MD}}$	decelerating (open-box enrichment)

External halo parameters for the models presented in Table 1. Table headers: (a) steepness parameter, (b) concentration parameter, (c) 'disc' radius (d) 'disc' mass, (e) total mass, and (f) Regime. In model 3, R_{hm} corresponds to the half-mass radius.

differ for $r < r_{\text{sc}}$; however, the discrepancy will be reconciled at larger radii, and thus one could talk of an 'asymptotically universal' property, in the sense defined by Salucci et al. (2007).

From the theoretical point of view, the importance of the analytic solution here presented resides in that it can be used to construct approximated superwind solutions when departures from spherical symmetry are important. In such a case, not just the parameters of the galaxy but also its morphology will determine both the superwind hydrodynamics and the fate of the ejected gas; specially, of metals. The connection of these features with the observed dispersion of the M-Z relationship will be discussed in a forthcoming work.

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